Property of Chiral Scalar and Axial-Vector Mesons in Heavy-Light Quark System

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We investigate the properties of new scalar X_B, X_D and axial-vector X_{B^*}, X_{D^*} mesons, which are recently predicted, in the covariant meson-classification scheme, to exist as the partners of ground-state pseudoscalar B, D and vector B^*, D^* mesons, respectively, realizing a linear representation of chiral symmetry for the light quark component in heavy-light-quark meson system. The mass splittings between the respective chiral partners are shown to be equal, and the decay widths of one pion emission of X_B, X_D, X_{B^*} and X_{D^*} become to have a common value, due to chiral and heavy quark symmetry.

§1. Introduction

The property of heavy-light quark meson systems is one of the important topics of high energy physics. Until recently the mass spectra and weak transitions of B and D meson systems have been analyzed mainly by using the heavy quark effective theory (HQET).[1] The HQET is based on heavy quark symmetry (HQS), which represent a spin-flavor symmetry of QCD in the heavy quark sector.

The QCD has another symmetry, chiral symmetry, in the light quark sector. Thus far for many years, the chiral symmetry had been applied in hadron physics mainly to describe the properties of π-meson octet through non-linear representation. Recently the existence of light σ-meson, $\sigma(500–600)$, is suggested strongly by the several groups,[2] including ours.[3] The observed properties of σ-meson are consistent[4] with the relation between mass and width predicted by the linear σ model (LσM),[5] and this fact seems to suggest that the linear representation of chiral symmetry is also realized[6] in nature.

In a previous paper[7] we have proposed a new covariant framework for meson level-classification and considered chiral transformation for the $q\bar{q}$ wave function (WF) directly in the framework of covariant oscillator quark model (COQM).[8] The σ is argued to be discriminated from the $(q\bar{q})$ $^3P_0$-state appearing in the non-relativistic quark model (NRQM), and to be classified into, so to speak, the "relativistic S-wave" state of the $(q\bar{q})$ system, which has never been appeared in NRQM. This type of new covariant state is called "chiralons." We assume the chiral symmetry for light quark component in the heavy-light quark systems is also effective, where the possible existence of this new type of mesons is suggested as chiral partners of the ground S-wave state pseudoscalar and vector mesons, respectively.

In this paper we investigate the properties of these new scalar and axial-vector

mesons by taking into account HQS for the heavy quark component, as well as chiral symmetry for the light quark component. Similar approach was given [9] from a viewpoint of the non-linear chiral Lagrangian, including only the ground-state psedoscalar and vector heavy-light quark mesons and the π-meson octet. In this work we adopt the linear representation, and make a chiral Lagrangian including also the new scalar and axial-vector mesons in the heavy-light B, D systems.

§2.   Constraints on One Pion Emission Process from HQS

For $n\bar{b} = (u\bar{b}, d\bar{b})$ systems the new scalar and axial-vector mesons are denoted as $X_B = {}^t(X_{B+}, X_{B^0})$ and $X_{B^*} = {}^t(X_{B^{*+}}, X_{B^{*0}})$. They are supposed to be the chiral partners of the pseudoscalar $B = {}^t(B^+, B^0)$ and vector $B^* = {}^t(B^{*+}, B^{*0})$, respectively. Similarly, for $n\bar{c} = {}^t(u\bar{c}, d\bar{c})$ system, the scalar $X_{\bar{D}} = {}^t(X_{\bar{D}^0}, X_{D-})$ and the axial-vector $X_{\bar{D}^*} = {}^t(X_{\bar{D}^{*0}}, X_{D^{*-}})$ are supposed to be the chiral partners of pseudoscalar $\bar{D} = {}^t(\bar{D}^0, D^-)$ and vector $\bar{D}^* = {}^t(\bar{D}^{*0}, D^{*-})$, respectively.

We consider general constraints from HQS on the processes of one pion emission, $X_B \to B\pi$, $X_{B^*} \to B^*\pi$, $X_D \to D\pi$ and $X_{D^*} \to D^*\pi$, which are expected to be the main decay modes of the relevant mesons.

The HQ spin symmetry relates B and X_B, respectively, to B^* and X_{B^*}, and D and X_D, respectively, to D^* and X_{D^*}, with the same velocity. Thus, the S-matrix elements of the relevant decay modes satisfy the relations

$$\langle \pi B(\mathbf{0})|U_I|X_B(\mathbf{0})\rangle = -i\langle \pi B^*(\mathbf{0}, \epsilon^{(0)})|U_I|X_{B^*}(\mathbf{0}, \epsilon^{(0)})\rangle,$$
$$\langle \pi D(\mathbf{0})|U_I|X_D(\mathbf{0})\rangle = -i\langle \pi D^*(\mathbf{0}, \epsilon^{(0)})|U_I|X_{D^*}(\mathbf{0}, \epsilon^{(0)})\rangle, \qquad (2\cdot1)$$

where U_I is translational operator of the time from $-\infty$ to $+\infty$, and $\mathbf{0}$ means the three velocity $\boldsymbol{v} = \mathbf{0}$. The longitudinally polarized states of B^*, D^* and X_{B^*}, X_{D^*} appear in RHS. When the z-direction is taken as a quantization axis, the above equations are derived by using the relation, $[U_I, S_Q^z] = 0$, $(2S_Q^z)^2 = 1$, $2S_Q^z|X_B(\mathbf{0})\rangle = |X_{B^*}(\mathbf{0}, \epsilon^{(0)})\rangle$ and $\langle B(\mathbf{0})|(2S_Q^z) = -i\langle B^*(\mathbf{0}, \epsilon^{(0)})|$ etc., where S_Q^z is the z-component of the heavy quark spin operator. (The factor $-i$ comes from the convention of the spinor WF [7] given in the appendix.) We assume that the mass difference between the chiral partners, $\Delta m_M = m_{X_M} - m_M$ ($M = B, D, B^*, D^*$), is sufficiently small compared to m_M, that is, $\Delta m_M \ll m_M \lesssim m_{X_M}$, but is larger than m_π. In this case, when we take the rest frame of initial X_M, the final heavy-light quark meson M also stays almost at rest, although the pion is emitted relativistically with energy Δm_M. Thus, in both sides of Eq.(2·1) the pion momentum $p_{\pi\mu}$ takes the same value, and the initial and final heavy-light quark mesons are at rest.

The HQ flavor symmetry relates the B, B^*, X_B and X_{B^*} mesons, respectively, to the D, D^*, X_D and X_{D^*} mesons with the same velocity. The HQ $\bar{b} \to \bar{c}$ flavor transformation leads to the following relations

$$\langle \pi B(\mathbf{0})|U_I|X_B(\mathbf{0})\rangle = \langle \pi D(\mathbf{0})|U_I|X_D(\mathbf{0})\rangle \equiv -\xi\frac{1}{(2\pi)^3 2}\sqrt{\frac{1}{(2\pi)^3 2E_\pi}}$$

$$\langle \pi B^*(\mathbf{0}, \epsilon^{(0)})|U_I|X_{B^*}(\mathbf{0}, \epsilon^{(0)})\rangle = \langle \pi D^*(\mathbf{0}, \epsilon^{(0)})|U_I|X_{D^*}(\mathbf{0}, \epsilon^{(0)})\rangle$$

$$\equiv -i\xi' \frac{1}{(2\pi)^3 2} \sqrt{\frac{1}{(2\pi)^3 2E_\pi}}, \qquad (2\cdot 2)$$

where a common factor $\frac{1}{(2\pi)^3 2}\sqrt{\frac{1}{(2\pi)^3 2E_\pi}}$ is introduced,[*] and the momentum conservation factor $i(2\pi)^4 \delta^{(4)}(P_{X_M} - P_M - p_\pi)$ is omitted in RHS. Due to Eq.(2·1), the ξ and ξ' satisfy

$$\xi = \xi'. \qquad (2\cdot 3)$$

Thus, all the relevant decay matrix elements are reduced to one universal amplitude ξ due to HQS.

The decay widths of $X_M = X_B, X_{B^*}, X_D, X_{D^*}$ are commonly calculated as

$$\Gamma_{X_M} = 3\frac{1}{2m_{X_M}} \frac{|\boldsymbol{p}|}{4\pi m_{X_M}} (\xi\sqrt{m_{X_M} m_M})^2 = 3\frac{\sqrt{(\Delta m_M)^2 - m_\pi^2}}{8\pi}\xi^2, \qquad (2\cdot 4)$$

where we use the approximation $m_{X_M} \approx m_M$, $|\boldsymbol{p}| = \sqrt{(\Delta m_M)^2 - m_\pi^2}$ is the pion CM momentum and the factor 3 comes from the isospin freedom of final $|\pi M\rangle$ state. Due to Eq.(2·4) the decay widths of the relevant processes are only dependent upon the correponding mass difference Δm_M.

§3. Linear Chiral Lagrangian for B and D Systems

We can simply make a linear chiral Lagrangian for the ground state B and D meson systems, if we know their transformation law. We shall use the following rules, which are derived by using the covariant spinor WF, [7] as explained in the appexdix.

The chiral transformation for the light u, d quark fields is given by

$$n = \begin{pmatrix} u \\ d \end{pmatrix} \rightarrow g_V \begin{pmatrix} u \\ d \end{pmatrix}, \quad g_A \begin{pmatrix} u \\ d \end{pmatrix}, \qquad (3\cdot 1)$$

where $g_V = \exp\{i\frac{\boldsymbol{\alpha}\cdot\boldsymbol{\tau} + \alpha_0 \tau^0}{2}\}$ and $g_A = \exp\{i\frac{\boldsymbol{\beta}\cdot\boldsymbol{\tau} + \beta_0 \tau^0}{2}\gamma_5\}$. Correspondingly, the chiral transformation for the light-heavy B and X_B meson fields is given by

$$(X_B + i\gamma_5 B) \rightarrow g_V(X_B + i\gamma_5 B), \ g_A(X_B + i\gamma_5 B),$$
$$(X_{B^*} + \gamma_5 B^*) \rightarrow g_V(X_{B^*} + \gamma_5 B^*), \ g_A(X_{B^*} + \gamma_5 B^*), \qquad (3\cdot 2)$$

which is also written in the form

$$(X_B - i\gamma_5 B) \rightarrow g_V(X_B - i\gamma_5 B), \ g_A^\dagger(X_B - i\gamma_5 B),$$
$$(X_{B^*} - \gamma_5 B^*) \rightarrow g_V(X_{B^*} - \gamma_5 B^*), \ g_A^\dagger(X_{B^*} - \gamma_5 B^*). \qquad (3\cdot 3)$$

[*] Here we use the ordinary normalization of the state $|B(\boldsymbol{0})\rangle \equiv a_{B(\boldsymbol{0})}^\dagger|0\rangle$ etc., where $[a_{B(\boldsymbol{p})}, a_{B(\boldsymbol{p'})}^\dagger] = \delta^{(3)}(\boldsymbol{p} - \boldsymbol{p'})$. Through heavy quark $\bar{b} \rightarrow \bar{c}$ transformation the state $|B(\boldsymbol{0})\rangle$ is transformed into $|\bar{D}(\boldsymbol{0})\rangle$, since one B meson is changed into one $\bar{D}$ meson. The amplitudes ξ and ξ', defined in Eq.(2·2), are related with the relevant Lorentz-invariant amplitudes, denoted respectively as $g_{X_B \rightarrow B}$, $g_{X_D \rightarrow D}$, $g_{X_{B^*} \rightarrow B^*}$ and $g_{X_{D^*} \rightarrow D^*}$, through the formula $g_{X_B \rightarrow B} = \sqrt{m_{X_B} m_B}\xi$, $g_{X_D \rightarrow D} = \sqrt{m_{X_D} m_D}\xi$, $g_{X_{B^*} \rightarrow B^*} = \sqrt{m_{X_{B^*}} m_{B^*}}\xi'$ and $g_{X_{D^*} \rightarrow D^*} = \sqrt{m_{X_{D^*}} m_{D^*}}\xi'$.

The chiral transformation for the charge conjugate $\bar{B} = (B^-, \bar{B}^0)$ and $X_{\bar{B}} = (X_{B^-}, X_{\bar{B}})$ meson fields is consistently given by

$$(X_{\bar{B}} + \bar{B}i\gamma_5) \rightarrow (X_{\bar{B}} + \bar{B}i\gamma_5)g_V^\dagger, \ (X_{\bar{B}} + \bar{B}i\gamma_5)g_A,$$
$$(X_{\bar{B}} - \bar{B}i\gamma_5) \rightarrow (X_{\bar{B}} - \bar{B}i\gamma_5)g_V^\dagger, \ (X_{\bar{B}} - \bar{B}i\gamma_5)g_A^\dagger. \tag{3.4}$$

$$(X_{\bar{B}^*} - \bar{B}^*\gamma_5) \rightarrow (X_{\bar{B}^*} - \bar{B}^*\gamma_5)g_V^\dagger, \ (X_{\bar{B}^*} - \bar{B}^*\gamma_5)g_A,$$
$$(X_{\bar{B}^*} + \bar{B}^*\gamma_5) \rightarrow (X_{\bar{B}^*} + \bar{B}^*\gamma_5)g_V^\dagger, \ (X_{\bar{B}^*} + \bar{B}^*\gamma_5)g_A^\dagger. \tag{3.5}$$

The chiral transformation rule of σ and $\boldsymbol{\pi}$ fields is given by

$$(\sigma + i\gamma_5\boldsymbol{\tau} \cdot \boldsymbol{\pi}) \rightarrow g_V(\sigma + i\gamma_5\boldsymbol{\tau} \cdot \boldsymbol{\pi})g_V^\dagger, \ g_A(\sigma + i\gamma_5\boldsymbol{\tau} \cdot \boldsymbol{\pi})g_A, \tag{3.6}$$
$$(\sigma - i\gamma_5\boldsymbol{\tau} \cdot \boldsymbol{\pi}) \rightarrow g_V(\sigma - i\gamma_5\boldsymbol{\tau} \cdot \boldsymbol{\pi})g_V^\dagger, \ g_A^\dagger(\sigma - i\gamma_5\boldsymbol{\tau} \cdot \boldsymbol{\pi})g_A^\dagger. \tag{3.7}$$

Then a chiral symmetric Lagrangian for B and X_B mesons is given by

$$\mathcal{L}_B = \mathcal{L}_B^{\text{free}} + \mathcal{L}_B^{\text{Yukawa}}, \tag{3.8}$$

where the free part $\mathcal{L}_B^{\text{free}}$ and the Yukawa coupling part $\mathcal{L}_B^{\text{Yukawa}}$ are given,[*] respectively, by

$$\mathcal{L}_B^{\text{free}} = -\frac{1}{4}\langle(\partial_\mu X_{\bar{B}} - \partial_\mu \bar{B}i\gamma_5)(\partial_\mu X_B + i\gamma_5\partial_\mu B)\rangle$$
$$- \frac{m_{B0}^2}{4}\langle(X_{\bar{B}} - \bar{B}i\gamma_5)(X_B + i\gamma_5 B)\rangle,$$
$$\mathcal{L}_B^{\text{Yukawa}} = -\frac{g_B}{4}\langle(X_{\bar{B}} - \bar{B}i\gamma_5)(\sigma + i\gamma_5\boldsymbol{\tau} \cdot \boldsymbol{\pi})(X_B - i\gamma_5 B)\rangle, \tag{3.9}$$

the notation $\langle \ \rangle$ denoting to take trace for spinor and light quark flavor indices. The B-meson and X_B meson have a common bare-mass m_{B0} due to chiral symmetry.

For the vector B_μ^* and axial-vector $X_{B_\mu^*}$ meson systems the chiral symmetric Lagrangian is given[**] by

$$\mathcal{L}_{B^*} = \mathcal{L}_{B^*}^{\text{free}} + \mathcal{L}_{B^*}^{\text{Yukawa}},$$
$$\mathcal{L}_{B^*}^{\text{free}} = -\frac{1}{8}\langle(F_{\mu\nu}^{X_{\bar{B}^*}} - F_{\mu\nu}^{\bar{B}^*}\gamma_5)(F_{\mu\nu}^{X_{B^*}} - \gamma_5 F_{\mu\nu}^{B^*})\rangle$$
$$- \frac{m_{B^*0}^2}{4}\langle(X_{\bar{B}^*\mu} - \bar{B}_\mu^*\gamma_5)(X_{B^*\mu} - \gamma_5 B_\mu^*)\rangle,$$
$$\mathcal{L}_{B^*}^{\text{Yukawa}} = -\frac{g_B^*}{4}\langle(X_{\bar{B}^*\mu} + \bar{B}_\mu^*\gamma_5)(\sigma + i\gamma_5\boldsymbol{\tau} \cdot \boldsymbol{\pi})(X_{B^*\mu} - \gamma_5 B_\mu^*)\rangle, \tag{3.10}$$

where $F_{\mu\nu}$ is the field strength tensor of the corresponding vector or axial vector mesons. The B^*-meson and X_{B^*}-meson have a common bare-mass m_{B^*0} in $\mathcal{L}_{B^*}^{\text{free}}$

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[*] The other chiral symmetric combinations, $\langle(\partial_\mu X_{\bar{B}} + \partial_\mu \bar{B}i\gamma_5)(\partial_\mu X_B - i\gamma_5\partial_\mu B)\rangle$, $\langle(X_{\bar{B}} + \bar{B}i\gamma_5)(X_B - i\gamma_5 B)\rangle$ and $\langle(X_{\bar{B}} + \bar{B}i\gamma_5)(\sigma - i\gamma_5\boldsymbol{\tau} \cdot \boldsymbol{\pi})(X_B + i\gamma_5 B)\rangle$ give the same contributions as the corresponding terms in Eq.(3.9), and thus omitted.

[**] The other chiral symmetric combinations, $\langle(F_{\mu\nu}^{X_{\bar{B}^*}} + F_{\mu\nu}^{\bar{B}^*}\gamma_5)(F_{\mu\nu}^{X_{B^*}} + \gamma_5 F_{\mu\nu}^{B^*})\rangle$, $\langle(X_{\bar{B}^*\mu} + \bar{B}_\mu^*\gamma_5)(X_{B^*\mu} + \gamma_5 B_\mu^*)\rangle$ and $\langle(X_{\bar{B}^*\mu} - \bar{B}_\mu^*\gamma_5)(\sigma - i\gamma_5\boldsymbol{\tau} \cdot \boldsymbol{\pi})(X_{B^*\mu} + \gamma_5 B_\mu^*)\rangle$, give the same contributions as the corresponding terms in Eq.(3.10), and thus omitted.

due to chiral symmetry. The spin-spin splitting is considered by m_{B0}^* being different from m_{B0} in Eq.(3·9).

Similarly, for D and D^* meson systems, the chiral Lagrangian is obtained from $\mathcal{L}_B$ and $\mathcal{L}_{B^*}$ simply by the substitution,

$$
\begin{aligned}
\mathcal{L}_D &= \mathcal{L}_B|_{(B,X_B,m_{B0},g_B)\to(\bar{D},X_{\bar{D}},m_{D0},g_D)} \\
\mathcal{L}_{D^*} &= \mathcal{L}_{B^*}|_{(B^*,X_{B^*},m_{B^*0},g_B^*)\to(\bar{D}^*,X_{\bar{D}^*},m_{D^*0},g_D^*)}.
\end{aligned} \tag{3·11}
$$

The total Lagrangian $\mathcal{L}^{\text{tot}}$ for ground-state B and D systems is given by

$$
\mathcal{L}^{\text{tot}} = \mathcal{L}_B + \mathcal{L}_D + \mathcal{L}_{B^*} + \mathcal{L}_{D^*} + \mathcal{L}_{L\sigma M}, \tag{3·12}
$$

where $\mathcal{L}_{L\sigma M}$ is the Lagrangian of SU(2) LσM,[4] describing the spontaneous chiral symmetry breaking. Apparently, Eq. (3·12) is possibly the most simple Lagrangian with interaction satisfying chiral symmetry.

Before going further, we consider the constraints on the present Lagrangian Eq.(3·12) required from HQS. The heavy quark spin symmetry and flavor symmetry predict, in this particular model, respectively, the relations among the relevant one-pion emission amplitudes,

$$
\begin{aligned}
& -g_B\sqrt{\frac{1}{(2\pi)^9}}\sqrt{\frac{1}{2m_{X_B}2m_B2E_\pi}} = -g_B^*\sqrt{\frac{1}{(2\pi)^9}}\sqrt{\frac{1}{2m_{X_{B^*}}2m_{B^*}2E_\pi}} \\
&= -g_D\sqrt{\frac{1}{(2\pi)^9}}\sqrt{\frac{1}{2m_{X_D}2m_D2E_\pi}} = -g_D^*\sqrt{\frac{1}{(2\pi)^9}}\sqrt{\frac{1}{2m_{X_{D^*}}2m_{D^*}2E_\pi}} \\
&\equiv -\eta\frac{1}{(2\pi)^32}\sqrt{\frac{1}{(2\pi)^32E_\pi}},
\end{aligned} \tag{3·13}
$$

where the universal amplitude η in the present scheme corresponds to the amplitude ξ defined in the general case (Eq.(2·3)).

§4.   Properties of Scalar and Axial-Vector Mesons

In the mechanism of spontaneous chiral symmetry breaking the σ acquires vacuum expectation value $\langle\sigma\rangle_0 \equiv \sigma_0 = f_\pi$, which is described by $\mathcal{L}_{L\sigma M}$ in Eq.(3·12). Through the $\mathcal{L}^{\text{Yukawa}}$ the mass degeneracies between scalar and pseudoscalar and between vector and axial-vector are broken as

$$
\begin{aligned}
m_{X_B}^2 &= m_{B0}^2 + g_B f_\pi, & m_B^2 &= m_{B0}^2 - g_B f_\pi, \\
m_{X_D}^2 &= m_{D0}^2 + g_D f_\pi, & m_D^2 &= m_{D0}^2 - g_D f_\pi, \\
m_{X_{B^*}}^2 &= m_{B^*0}^2 + g_B^* f_\pi, & m_{B^*}^2 &= m_{B^*0}^2 - g_B^* f_\pi, \\
m_{X_{D^*}}^2 &= m_{D^*0}^2 + g_D^* f_\pi, & m_{D^*}^2 &= m_{D^*0}^2 - g_D^* f_\pi,
\end{aligned} \tag{4·1}
$$

which leads to

$$
\Delta m_B = \Delta m_{B^*} = \Delta m_D = \Delta m_{D^*} \equiv \Delta m, \tag{4·2}
$$

because of Eq.(3·13), where the approximations $m_{X_B} \approx m_B \approx m_{B0}$ etc. are used. Due to Eq.(2·4), Eq.(4·2) leads to the universality of decay widths,

$$\Gamma_{X_B} = \Gamma_{X_D} = \Gamma_{X_{B^*}} = \Gamma_{X_{D^*}} = \Gamma(\Delta m). \tag{4·3}$$

In this model the Δm is given by

$$\Delta m = \eta f_\pi, \tag{4·4}$$

and thus, the $\Gamma(\Delta m)$ is given by Eq.(2·4)

$$\Gamma(\Delta m) = = 3\frac{\sqrt{(\Delta m)^2 - m_\pi^2}}{8\pi}\xi^2, \text{ with } \xi = \eta = \Delta m/f_\pi. \tag{4·5}$$

The numerical value of $\Gamma(\Delta m)$ (with $f_\pi = 93\text{MeV}$) is depicted by solid line in Fig. 1.

The universality of mass splittings between the relevant chiral partners and of the decay widths, which was derived by using the Yukawa coupling in the most simple form, Eqs.(3·9) and (3·10), is shown to be rather generally valid from the following considerations. The mass-splittings and decay widths can be described more generally in the present scheme of spontaneous chiral symmetry breaking by the following Lagrangian,

$$\mathcal{L}_n^{\text{Yukawa}} = -\sum_{M=B,\bar{D}} \frac{g_{M,n}}{4}\langle(X_{\bar{M}} - \bar{M}i\gamma_5)(\sigma + i\gamma_5\boldsymbol{\tau} \cdot \boldsymbol{\pi})(\sigma^2 + \boldsymbol{\pi}^2)^n(X_M - i\gamma_5 M)\rangle$$

$$-\sum_{M=B^*,\bar{D}^*} \frac{g_{M,n}}{4}\langle(X_{\bar{M}\mu} + \bar{M}_\mu\gamma_5)(\sigma + i\gamma_5\boldsymbol{\tau} \cdot \boldsymbol{\pi})(\sigma^2 + \boldsymbol{\pi}^2)^n(X_{M\mu} - \gamma_5 M_\mu)\rangle \tag{4·6}$$

$$\mathcal{L}_n^{\text{Z}} = \sum_{M=B,\bar{D}} \frac{h_{M,n}}{4}\langle\partial_\mu(X_{\bar{M}} - \bar{M}i\gamma_5)(\sigma + i\gamma_5\boldsymbol{\tau} \cdot \boldsymbol{\pi})(\sigma^2 + \boldsymbol{\pi}^2)^n\partial_\mu(X_M - i\gamma_5 M)\rangle$$

$$+\sum_{M=B^*,\bar{D}^*} \frac{h_{M,n}}{4}\langle\partial_\mu(X_{\bar{M}\nu} + \bar{M}_\nu\gamma_5)(\sigma + i\gamma_5\boldsymbol{\tau} \cdot \boldsymbol{\pi})(\sigma^2 + \boldsymbol{\pi}^2)^n\partial_\mu(X_{M\nu} - \gamma_5 M_\nu)\rangle \tag{4·7}$$

with $n = 0, 1, 2, \dots$. When the σ takes VEV, $\sigma_0 = f_\pi$, the effects from the terms with $n \geq 2$ on the mass splittings and on the one pion emission amplitude can be included in the $n = 0$ term through the redefinition,

$$g_{M,0} \to g_M \equiv \sum_{n=0,1,2,\dots} g_{M,n}f_\pi^{2n} \text{ and } h_{M,0} \to h_M \equiv \sum_{n=0,1,2,\dots} h_{M,n}f_\pi^{2n}. \tag{4·8}$$

The modified $n = 0$ term of Eq.(4·6) is the simplest Yukawa Lagrangian, given in Eqs. (3·9) and (3·10), leading to the mass formula, Eq.(4·1), and to the one pion emission amplitude, Eq.(3·13). The modified $n = 0$ term of Eq.(4·7) leads to the field renormalization,

$$\sqrt{1 - h_M f_\pi}X_M \equiv X_M^{(r)}, \quad \sqrt{1 + h_M f_\pi}M \equiv M^{(r)}, \quad (M = B, D, B^*, D^*), \tag{4·9}$$

which gives, in the case of $h_M f_\pi \ll 1$, the same mass formula as Eq.(4·1) with g_M replaced by $h_M m_{M0}^2$. This term also gives the one pion emission amplitude

$$h_M P_{X_M} \cdot P_M \sqrt{\frac{1}{(2\pi)^9}}\sqrt{\frac{1}{2E_{X_M}2E_M 2E_\pi}} \approx -h_M m_{X_M} m_M \sqrt{\frac{1}{(2\pi)^9}}\sqrt{\frac{1}{2m_{X_M}2m_M 2E_\pi}}$$
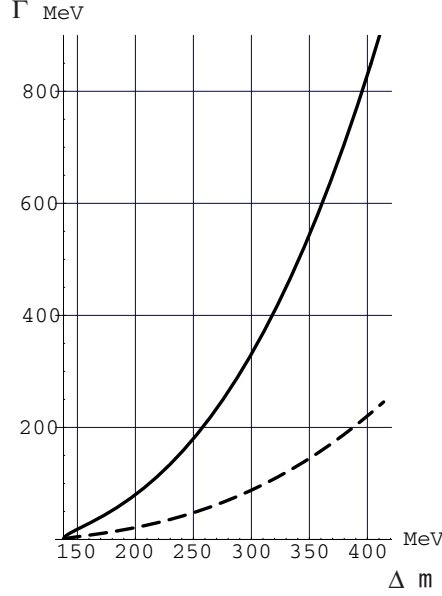
Fig. 1. Δm-dependence of the universal one-pion-emission decay width $\Gamma(\Delta m)$ of X_B, X_D, X_{B^*} and X_{D^*} mesons: The $\Delta m = m_{X_M} - m_M$ ($M = B, D, B^*, D^*$) and $\Gamma(\Delta m)$ take a common value due to HQS and chiral symmetry. The solid line represents the prediction by using the Lagrangian in Eqs.(3·9) and (3·10). The dashed line represents the prediction in the case of including the correction of Eq.(4·11) with the $k = 1/m_q^{\mathrm{constit}} = 2/m_\rho$.

$$\approx -h_M m_{M0}^2 \sqrt{\frac{1}{(2\pi)^9}} \sqrt{\frac{1}{2m_{X_M} 2m_M 2E_\pi}}, \tag{4·10}$$

which is given from the LHS of Eq.(3·13) by replacing g_M with $h_M m_{M0}^2$.

Thus, the effects from Eqs.(4·6) and (4·7) can be described by the simplest Yukawa Lagrangian in Eqs. (3·9) and (3·10). Thus, its result of the universality of mass-splittings Eq.(4·2), and of the universality of decay widths Eq.(4·3) is kept to be valid even in the case of taking the general contributions from Eqs. (4·6) and (4·7) into account. However, the value of universal $\Gamma(\Delta m)$ itself generally becomes different from the original one, Eq.(4·5), by considering the effect from the following type of Lagrangian,

$$\mathcal{L}^{(d)} = \sum_{M=B,\bar{D}} \frac{k_M}{4} \langle \{(X_M - i\gamma_5 M) \overset{\leftrightarrow}{\partial_\mu} (X_{\bar{M}} - \bar{M}i\gamma_5)\} \partial_\mu(\sigma + i\gamma_5 \boldsymbol{\tau} \cdot \boldsymbol{\pi}) \rangle$$

$$+ \sum_{M=B^*,\bar{D}^*} \frac{k_M}{4} \langle \{(X_{M\nu} - \gamma_5 M_\nu) \overset{\leftrightarrow}{\partial_\mu} (X_{\bar{M}\nu} + \bar{M}_\nu \gamma_5)\} \partial_\mu(\sigma + i\gamma_5 \boldsymbol{\tau} \cdot \boldsymbol{\pi}) \rangle, \tag{4·11}$$

where $k_B = k_B^* = k_D = k_D^* \equiv k$ due to HQS. By taking this Lagrangian into account, the Lorentz invariant amplitude of one pion emission decay $-g_M(= -\eta\sqrt{m_{X_M} m_M} \approx -\eta\, m_{X_M} = -\Delta m\, m_{X_M}/f_\pi$ due to Eq.(3·13)) is changed as

$$-g_M \rightarrow -g_M - k\, p_\pi \cdot (P_M + P_{X_M}) \approx -g_M + 2k\Delta m\, m_{X_M}$$

$$= -\Delta m \; m_{X_M}(\frac{1}{f_\pi} - 2k), \qquad (4 \cdot 12)$$

and the $\Gamma(\Delta m)$ is still given by Eq.(4·5) with the new ξ

$$\xi = \Delta m(1/f_\pi - 2k). \qquad (4 \cdot 13)$$

The order of the magnitude of k is inverse momentum of the light quark. The $\Gamma(\Delta m)$ is given in the case of $k = 1/m_q^{\text{constit}}($, where the constituent quark mass $m_q^{\text{constit}} \approx m_\rho/2)$, as an example, by dashed line in Fig. 1. As can be seen in this figure the absolute value of Γ is largely dependent upon the value of k, and cannot be predicted in the present model.

Finally we consider the possible effects of mixing between pseudo-scalar and axial-vector and between scalar and vector mesons. The corresponding Lagrangian is given by

$$\mathcal{L}^{\text{mix}} = \sum_{M=B,\bar{D}} \frac{l_M}{4} \langle (X_{\bar{M}^*\mu} + \bar{M}^*_\mu \gamma_5) i\gamma_5 (\partial_\mu X_M + i\gamma_5 \partial_\mu M)$$

$$+ (\partial_\mu X_{\bar{M}} + \partial_\mu \bar{M} i\gamma_5) i\gamma_5 (X_{M^*\mu} - \gamma_5 M^*_\mu) \rangle$$

$$= \sum_M l_M [-(X_{\bar{M}^*\mu}\partial_\mu M + \partial_\mu \bar{M} X_{\bar{M}^*\mu}) + i(\bar{M}^*_\mu \partial_\mu X_M - \partial_\mu X_{\bar{M}} M^*_\mu)], \qquad (4 \cdot 14)$$

where $l_B = l_D \equiv l$ due to HQS, and the magnitude of this l is of the order of light quark momentum $\sim m_q^{\text{constit}}$. The equation (4·14) leads to the redefinition of the vector and axial-vector meson fields,

$$\boldsymbol{X}_{\bar{M}^*\mu} \equiv X_{\bar{M}^*\mu} + \frac{l}{m^2_{X_{M^*}}}\partial_\mu M, \quad \boldsymbol{X}_{M^*\mu} \equiv X_{M^*\mu} + \frac{l}{m^2_{X_{M^*}}}\partial_\mu M,$$

$$\bar{\boldsymbol{M}}^*_\mu \equiv \bar{M}^*_\mu + \frac{il}{m^2_{M^*}}\partial_\mu X_{\bar{M}}, \quad \boldsymbol{M}^*_\mu \equiv M^*_\mu - \frac{il}{m^2_{M^*}}\partial_\mu X_M. \qquad (4 \cdot 15)$$

Correspondingly, the field renormalization for the scalar and pseudoscalar mesons is necessary.

$$\sqrt{1 - \frac{l^2}{m^2_{X_{M^*}}}}M \to M, \quad \sqrt{1 - \frac{l^2}{m^2_{M^*}}}X_M \to X_M. \qquad (4 \cdot 16)$$

The field renormalization changes the masses of scalar and pseudoscalar mesons commonly as $m^2_M \to m^2_M(1 + l^2/m^2_M)$ and $m^2_{X_M} \to m^2_{X_M}(1 + l^2/m^2_{X_M})$. This correction factor is of order $(m_q/m_Q)^2$, and thus expected to be small. Furthermore, it does not affect the mass splitting Δm_M. Thus, it is not necessary to take these mixing effects into account.

§5. Concluding Remarks

We have investigated the properties of new scalar X_B, X_D and axial vector X_{B^*}, X_{D^*} mesons, which are chiral partners of the ground state pseudoscalar B, D and vector B^*, D^* mesons, respectively.

Through the mechanism of spontaneous chiral symmetry breaking the masses of chiral multiplets split universally, and the decay widths for one pion emission of X_B, X_D, X_{B^*} and X_{D^*} generally take a common value in the present framework, due to chiral and heavy quark symmetry.

Possible existence of these new mesons and the validity of above predicted property should be investigated from experimental data. Recently the spectra of $D^*\pi$ and $B\pi$ invariant mass from weak Z^0-boson decay were analyzed, and some evidences for possible existence of the new axial-vector X_{D^*} and scalar X_B are suggested, respectively. The preliminary values of masses are $m_{X_{D^*}} \approx 2300\text{MeV}$ [10] and $m_{X_B} \sim 5520\text{MeV}$, [11] corresponding to $\Delta m_{D^*} \approx 290\text{MeV}$ and $\Delta m_B \sim 240\text{MeV}$, respectively. However, in order to derive any definite conclusion, the data with higher statistics are necessary.

Acknowledgements

Appendix A
—— *Derivation of Chiral and HQ Spin Transformation Rule from Covariant Spinor WF* ——

In this appendix we derive the chiral and HQ spin transformation rules for the heavy-light quark mesons which are, *up to the phase factor*, derived from the covariant bi-spinor WF given in our previous paper. [7] The (annihilation part of) WF for the light-heavy $q\bar{Q}$ mesons, with four velocity v_μ are

$$(U_B^{(+)}, D_{X_B}^{(+)}, U_{B^*(s)}^{(+)}, D_{X_{B^*}^{(s)}}^{(+)}) = \frac{1}{2\sqrt{2}}(i\gamma_5, 1, i\epsilon^{(s)} \cdot \gamma, i\gamma_5\epsilon^{(s)} \cdot \gamma)(1 + iv \cdot \gamma), \quad \text{(A·1)}$$

respectively. (The index $(+)$ is omitted in the following.) Corresponging to the "$i\gamma_5$"-transformation for the light-quark component WF, $u(\boldsymbol{p}, s) \to u'(\boldsymbol{p}, s) = i\gamma_5 u(\boldsymbol{p}, s)$, the above bi-spinor meson WF are transformed as

$$U_B \to i\gamma_5 U_B = -D_{X_B}, \quad D_{X_B} \to i\gamma_5 D_{X_B} = U_B,$$
$$U_{B^*(s)} \to i\gamma_5 U_{B^*(s)} = iD_{X_{B^*}^{(s)}}, \quad D_{X_{B^*}^{(s)}} \to i\gamma_5 D_{X_{B^*}^{(s)}} = iU_{B^*(+)}, \quad \text{(A·2)}$$

which are represented by the χ field operation as

$$\chi a_{B(\boldsymbol{p})}^\dagger |0\rangle = -a_{X_B(\boldsymbol{p})}^\dagger |0\rangle, \quad \chi a_{X_B(\boldsymbol{p})}^\dagger |0\rangle = a_{B(\boldsymbol{p})}^\dagger |0\rangle,$$
$$\chi a_{B^*(\boldsymbol{p},\epsilon^{(s)})}^\dagger |0\rangle = i a_{X_{B^*}(\boldsymbol{p},\epsilon^{(s)})}^\dagger |0\rangle, \quad \chi a_{X_{B^*}(\boldsymbol{p},\epsilon^{(s)})}^\dagger |0\rangle = i a_{B^*(\boldsymbol{p},\epsilon^{(s)})}^\dagger |0\rangle. \quad \text{(A·3)}$$

In terms of the local meson field operator, Eq.(A·3) is represented by

$$\chi(\bar{B}, X_{\bar{B}}, \bar{B}^*, X_{\bar{B}^*})\chi^\dagger = (-X_{\bar{B}}, \bar{B}, iX_{\bar{B}^*}, i\bar{B}^*). \quad \text{(A·4)}$$

Through its hermitian conjugation (which converts the charge of the meson), Eq.(A·4) leads to

$$\chi(B, X_B, B^*, X_{B^*})\chi^\dagger = (-X_B, B, -iX_{B^*}, -iB^*), \tag{A·5}$$

which is equivalent to

$$\chi^\dagger(B, X_B, B^*, X_{B^*})\chi = (X_B, -B, iX_{B^*}, iB^*). \tag{A·6}$$

The $\chi^\dagger$ field operation corresponds to the "$-i\gamma_5$" transformation for the quark component WF, $u(\boldsymbol{p}, s) \to u'(\boldsymbol{p}, s) = -i\gamma_5 u(\boldsymbol{p}, s)$. The equation (A·6) is also derived in a different way. The WF are given by

$$\langle 0|\Phi\, a^\dagger_{B(\boldsymbol{p})}|0\rangle, \;\; \langle 0|\Phi\, a^\dagger_{X_B(\boldsymbol{p})}|0\rangle, \;\; \langle 0|\Phi\, a^\dagger_{B^{*(s)}(\boldsymbol{p})}|0\rangle, \;\; \langle 0|\Phi\, a^\dagger_{X^{(s)}_{B^*}(\boldsymbol{p})}|0\rangle, \tag{A·7}$$

where

$$\Phi = \int d\boldsymbol{p}'(a_{B(\boldsymbol{p}')}U_B + a_{X_B(\boldsymbol{p}')}D_{X_B}$$
$$+ \sum_s (a_{B^*(\boldsymbol{p}',\epsilon^{(s)})}U_{B^{*(s)}} + a_{X_{B^*}(\boldsymbol{p}',\epsilon^{(s)})}D_{X^{(s)}_{B^*}})). \tag{A·8}$$

Under the $i\gamma_5$-transformation, $\langle 0|\Phi a^\dagger_{B(\boldsymbol{p})}|0\rangle \to i\gamma_5\langle 0|\Phi a^\dagger_{B(\boldsymbol{p})}|0\rangle$, which corresponds to $\langle 0|\Phi a^\dagger_{B(\boldsymbol{p})}|0\rangle \to \langle 0|\Phi\chi a^\dagger_{B(\boldsymbol{p})}|0\rangle = \langle 0|\chi\, \chi^\dagger\Phi\chi\, a^\dagger_{B(\boldsymbol{p})}|0\rangle$. Thus, the transformation rule is given by $\chi^\dagger\Phi\chi = i\gamma_5\Phi$, which is equivalent to

$$\chi^\dagger(Bi\gamma_5 + X_B 1 + B^* i\epsilon \cdot \gamma + X_{B^*}i\gamma_5\epsilon \cdot \gamma)\chi$$
$$= i\gamma_5(Bi\gamma_5 + X_B 1 + B^* i\epsilon \cdot \gamma + X_{B^*}i\gamma_5\epsilon \cdot \gamma). \tag{A·9}$$

The equation (A·9) is more simplified as

$$\chi^\dagger(Bi\gamma_5 + X_B 1 + B^*\gamma_5 + X_{B^*})\chi = i\gamma_5(Bi\gamma_5 + X_B 1 + B^*\gamma_5 + X_{B^*}), \tag{A·10}$$

which is equivalent to Eq.(A·6).

The $\chi^\dagger$ transformation rule for charge conjugated $b\bar{n}$ systems is given from the Pauli conjugate, $\Phi \to \gamma_4\Phi^\dagger\gamma_4$, of Eq.(A·10) as

$$\chi^\dagger(\bar{B}i\gamma_5 + X_{\bar{B}} 1 - \bar{B}^*\gamma_5 + X_{\bar{B}^*})\chi = (\bar{B}i\gamma_5 + X_{\bar{B}} 1 - \bar{B}^*\gamma_5 + X_{\bar{B}^*})i\gamma_5. \tag{A·11}$$

By using Eqs.(A·10) and (A·11) we can derive the chiral transformation rules given in §3(, which are the $\chi^\dagger$ transformation,) of the text.

Next we consider the operation of S^z_Q on our bi-spinor WF. We take the meson rest frame and the z-direction as the spin quantization axis. Representing the spin configurations as $|s^z_q, S^z_{\bar{Q}}\rangle$ ($s^z_q(S^z_{\bar{Q}})$ denoting the third component of light quark (heavy anti-quark) spin), the bi-spinor WF, $U_B, U_{B^{*(0)}}, D_{X_B}$ and $D_{X_{B^{*(0)}}}$, are given symbolically by

$$U_B = \frac{1}{i\sqrt{2}}(|+-\rangle - |-+\rangle), \quad U_{B^{*(0)}} = \frac{1}{\sqrt{2}}|+-\rangle + |-+\rangle,$$
$$D_{X_B} = \frac{1}{\sqrt{2}}|+-\rangle - |-+\rangle, \quad D_{X^{(0)}_{B^*}} = -\frac{1}{\sqrt{2}}(|+-\rangle + |-+\rangle), \tag{A·12}$$

where in U-type bi-spinor WF both the light-quark and heavy-antiquark spinors are in positive energy state, but in D-type bi-spinor WF the light-quark spinor is in "negative energy" state. The operation of S_Q^z on WF is the multiplication of $-\sigma_3^*$ to the $\bar{Q}$-index of WF from left side (or equivalent to the multiplication of $-\sigma_3^\dagger = -\sigma_3$ from right side).

$$U_{B(\mathbf{0})}(-\sigma_3) = iU_{B^{*(0)}(\mathbf{0})}, \quad U_{B^{*(0)}(\mathbf{0})}(-\sigma_3) = -iU_{B(\mathbf{0})},$$
$$D_{X_B(\mathbf{0})}(-\sigma_3) = D_{X_{B^*}^{(0)}(\mathbf{0})}, \quad D_{X_{B^*}^{(0)}(\mathbf{0})}(-\sigma_3) = D_{X_B(\mathbf{0})}. \tag{A·13}$$

The equation (A·13) corresponds to the following relation for meson creation operator

$$[2S_Q^z, a_{B(\mathbf{0})}^\dagger] = ia_{B^{*(0)}(\mathbf{0})}^\dagger, \quad [2S_Q^z, a_{B^{*(0)}(\mathbf{0})}^\dagger] = -ia_{B(\mathbf{0})}^\dagger,$$
$$[2S_Q^z, a_{X_B(\mathbf{0})}^\dagger] = a_{X_{B^*}^{(0)}(\mathbf{0})}^\dagger, \quad [2S_Q^z, a_{X_{B^*}^{(0)}(\mathbf{0})}^\dagger] = a_{X_B(\mathbf{0})}^\dagger. \tag{A·14}$$

The hermitian conjugate of the above equations leads to

$$[2S_Q^z, a_{B(\mathbf{0})}] = ia_{B^{*(0)}(\mathbf{0})}, \quad [2S_Q^z, a_{B^{*(0)}(\mathbf{0})}] = -ia_{B(\mathbf{0})},$$
$$[2S_Q^z, a_{X_B(\mathbf{0})}] = -a_{X_{B^*}^{(0)}(\mathbf{0})}, \quad [2S_Q^z, a_{X_{B^*}^{(0)}(\mathbf{0})}] = -a_{X_B(\mathbf{0})}. \tag{A·15}$$

The equations (A·14) and (A·15) are represented by using meson field operator symbolically as

$$[2S_Q^z, \bar{B}] = i\bar{B}^{*(0)}, \quad [2S_Q^z, \bar{B}^{*(0)}] = -i\bar{B},$$
$$[2S_Q^z, X_{\bar{B}}] = X_{\bar{B}^*}^{(0)}, \quad [2S_Q^z, X_{\bar{B}^*}^{(0)}] = X_{\bar{B}}; \quad \text{and}$$
$$[2S_Q^z, B] = iB^{*(0)}, \quad [2S_Q^z, B^{*(0)}] = -iB,$$
$$[2S_Q^z, X_B] = -X_{B^*}^{(0)}, \quad [2S_Q^z, X_{B^*}^{(0)}] = -X_B, \tag{A·16}$$

where only the components with velocity $\boldsymbol{v} = \mathbf{0}$ are selected in all fields. The above equations are useful to make the Lagrangian satisfying the heavy quark spin symmetry.

References

1) N. Isgur and M. B. Wise, Phys. Lett. **B232** (1990), 113; **B237** (1990), 527; Nucl. Phys. **B348** (1991), 276.
   A. F. Falk, H. Georgi, B. Grinstein and M. B. Wise, Nucl. Phys. **B343** (1991), 1.
   H. Georgi, Phys. Lett. **B240** (1990), 447; Nucl. Phys. **B348** (1991), 293.
   B. Grinstein, Nucl. Phys. **339** (1990), 253.
   M. Neubert, Phys. Lett. B264 (1991), 455; Phys. Rep. **245** (1994), 259.
2) N. N. Achasov and G. N. Schestakov, Phys. Rev. **D49** (1994), 5779.
   R. Kaminski, L. Lesniak and J. -P. Maillet, Phys. Rev. **D50** (1994), 3145.
   N. A. Tornqvist, Z. Phys. **C68** (1995), 647.
   N. A. Tornqvist and M. Roos, Phys. Rev. Lett. **76** (1996), 1575.
   M. Harada, F. Sannino and J. Schechter, Phys. Rev. **D54** (1996), 1991.
3) S. Ishida, M. Ishida, H. Takahashi, T. Ishida, K. Takamatsu and T. Tsuru, Prog. Theor. Phys. **95** (1996), 745; **98** (1997), 1005.

4) M. Ishida, Prog. Theor. Phys. **96** (1996), 853.
M. Ishida *in proc of Workshop on Hadron Spectroscopy, Frascati, 1999,* (WHS 99), ed. by
T. Bressani, A. Feliciello and A. Filippi, Frascati Physics Series XV (1999), p.85.
M. Ishida,*in proc of Possible Existence of σ-Meson and Its Implications to Hadron
Physics, YITP Kyoto, 2000,* (σ-Workshop), Soryushiron Kenkyu **102** (2001), No.5; KEK-
proceedings/2000-4.
See also, N. A. Tornqvist Phys.Lett.**B426** (1998), 105.
M. Ishida, Prog. Theor. Phys. **101** (1999), 661.
5) J. Schwinger, Ann. Phys. **2** (1957), 407.
M. Gell-Mann and M. Levi, Nuovo Cim. **16** (1960), 705.
6) R. Delbourgo and M. D. Scadron, Phys. Rev. Lett. **48** (1982), 765.
T. Hatsuda and T. Kunihiro, Prog. Theor. Phys. **74** (1985), 765; Phys. Rep. **247** (1994),
221.
7) S. Ishida, M. Ishida and T. Maeda, Prog. Theor. Phys. **104** (2000), 785.
S. Ishida, M. Ishida and T. Maeda, *in proc. of 8th Int. Conf. on Hadron Spectroscopy,
Beijing, China, 24-28 Aug. 1999,* (Hadron99); hep-ph/0002095.
S. Ishida and M. Ishida, in proc. of σ-Workshop in Kyoto.
8) S. Ishida, Prog. Theor. Phys. **45** (1971), 1570, 1905.
R. P. Feynmann, M. Kislinger and F. Ravndal, Phys. Rev. **D3** (1971), 2706.
H. Yukawa, Phys. Rev. **77** (1950), 219; **91** (1953), 415, 416.
T. Takabayashi, Nuovo Cim. **33** (1964), 668.
I. Sogami, Prog. Theor. Phys. **41** (1969), 1352.
K. Fujimura, T. Kobayashi and M. Namiki, Prog. Theor. Phys. **43** (1970), 73.
Y. S. Kim and M. E. Noz, Phys. Rev. **D8** (1973), 3251.
S. Ishida, M. Y. Ishida and M. Oda, Prog. Theor. Phys. **93** (1995), 939.
The COQM has a long history of development since the bilocal theory by H. Yukawa. On
the preceding works see, the references given in the above last paper.
9) M. B. Wise, Phys. Rev. **D45** (1992), 2188.
G. Burdman and J. Donoghue, Phys. Lett. **B280** (1992), 287.
T. M. Yan, H. Y. Cheng, C. Y. Cheung, G. L. Lin, Y. C. Lin and H. L. Yu, Phys. Rev.
D46 (1992), 1148.
10) H. Tonooka, master thesis, Nihon university, 2001.
11) T. Maeda, master thesis, Nihon university, 2000.